

$Z^0-Z^{0'}$ and ordinary-exotic fermion mixings revisited using some special observables

Gautam Bhattacharyya ^{*)}

*Theory Division, CERN,
CH 1211, Geneva 23,
SWITZERLAND*

ABSTRACT

We investigate the mixing of an extra $Z^{0'}$ with the standard Z^0 and mixings of exotic fermions with their standard counterparts through some precisely measured electroweak observables. These observables are geared to search for physics beyond the Standard Model. We observe that although most of such mixings are severely constrained by the recent LEP data, some of the mixing angles could still be rather large, awaiting future tests.

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^{*)} gautam@cernvm.cern.ch

1 Introduction

Precision measurements on and around the Z peak at LEP have already put severe constraints on the possible mixing effects of the standard Z^0 with one extra neutral vector boson $Z^{0'}$ and the mixings of sequential (ordinary) fermions with their exotic counterparts. These additional fermions and neutral bosons may stem from many a theoretical extension of the Standard Model (SM), the most popular ones being the superstring-inspired E_6 [1] and the left-right symmetric (LRS) model [2] which contain the SM gauge group within their group structures. During each phase of the LEP run these analyses were carried out [3–12]. Bounds have been put on Z^0 – $Z^{0'}$ mixing angles in refs. [3–6] and on ordinary-exotic mixing angles in refs. [7–9] (on both in some refs.), by comparing the modified theoretical predictions of $e^+e^- \rightarrow Z \rightarrow f\bar{f}$ cross sections and forward-backward asymmetries on and around the Z peak with the corresponding experimental results. Variation of those bounds with m_t , M_H and α_S have also been studied. The authors of refs. [10–12] have introduced a special set of observables, which are specifically geared to study these extensions beyond the (SM) with a rather focused attention. These new observables are constructed by combining the standard ones, which are directly probed by experiments, mainly at LEP. One of the most interesting features of these observables is that they can disentangle the two different sources of quadratic m_t dependences in neutral current processes, appearing through (i) $\Delta\rho \equiv \rho - 1$ (originating from two-point W and Z vacuum polarizations) and (ii) ΔV_b^t (arising out of the $Zb\bar{b}$ vertex correction). Refs. [11, 12] deal basically with extensions motivated from E_6 . In this paper we intend to impose bounds on more general kinds of mixing effects, updating existing bounds in some cases, using the latest LEP data [13]. As tools for our studies we use some currently fashionable observables which enjoy measurements of unprecedented precision at LEP.

2 The observables and their SM expressions

Before getting down to the business of new physics we first recall the different LEP observables that interest us and their SM expressions. The partial width of Z to any particular standard fermion flavour i is given by

$$\Gamma_i = N_c^i \frac{G_\mu M_Z^3}{6\pi\sqrt{2}} (v_i^2 + a_i^2), \quad (1)$$

where

$$\begin{aligned} N_c^i &= 1 + \frac{3\alpha}{4\pi} Q_i^2, & (i = \text{lepton}), \\ &= 3 \left(1 + \frac{3\alpha}{4\pi} Q_i^2 \right) \left[1 + \frac{\alpha_S(M_Z)}{\pi} + \dots \right], & (i = \text{quark}). \end{aligned} \quad (2)$$

The vector (v_i) and axial vector (a_i) couplings in the SM are given by

$$\begin{aligned} v_i^{\text{SM}} &= \sqrt{\rho} \left[t_3^i - 2Q_i \sin^2 \bar{\theta}_W \right] \\ a_i^{\text{SM}} &= \sqrt{\rho} t_3^i \end{aligned} \quad (3)$$

where $\sqrt{\rho}$ represents a non-trivial wave function renormalization of the on-shell Z . In the framework of the SM ρ is unity at the tree level and it experiences a quantum correction ($\Delta\rho_t$) mainly due to the t - b mass splitting. To a very good approximation

$$\Delta\rho_t \simeq \frac{3G_\mu m_t^2}{8\pi^2 \sqrt{2}}. \quad (4)$$

After improved Born approximation $\sin^2 \bar{\theta}_W$ is different from S_0^2 ($\simeq 0.23$), and for all practical purposes, $\sin^2 \bar{\theta}_W \simeq S_0^2 - \frac{3}{8} \Delta\rho_t$. The couplings of Z to b quarks require a further correction, which arises from the t -mediated triangle loop in the above vertex. Since the t quark couples in full strength with the b quark through W exchange, this non-universal correction is important only for the external b lines. The modified couplings are given by

$$\begin{aligned} v_b &= v_d - \frac{19}{60} \Delta V_b^t \\ a_b &= a_d - \frac{19}{60} \Delta V_b^t, \end{aligned} \quad (5)$$

where

$$\Delta V_b^t = -\frac{20\alpha}{19\pi} \left(\frac{m_t^2}{M_Z^2} + \frac{13}{6} \ln \frac{m_t^2}{M_Z^2} \right). \quad (6)$$

We now introduce four observables which are sensitive to many a different kind of new physics, including the ones we are probing for. We also specify the virtues of these observables. The first one ($\tilde{\gamma}_e$), which is called the reduced electronic width [11], is given by

$$\tilde{\gamma}_e = \gamma_e - \frac{2}{3} \xi \quad (7)$$

where

$$\gamma_e = \frac{9 \Gamma_e}{\alpha(M_Z) M_Z}; \quad \xi = \frac{M_W^2}{M_Z^2 C_0^2}; \quad C_0^2 = 1 - S_0^2. \quad (8)$$

In the SM:

$$\tilde{\gamma}_e^{\text{SM}} = \frac{1}{3} + \frac{2v_0}{3}, \quad \text{where } v_0 = 1 - 4S_0^2. \quad (9)$$

Notice that $\tilde{\gamma}_e$ is free from uncertainties originating from m_t or M_H .

The second one [11] $R_{\tau e} = \frac{\Gamma_\tau}{\Gamma_e}$ is a measure of any violation of lepton universality on the Z peak. In the SM (neglecting the τ mass), it is given by

$$R_{\tau e}^{\text{SM}} = 1. \quad (10)$$

The third observable $R_b = \Gamma_b/\Gamma_{had}$ has recently received wide attention from many quarters and its experimental measurement also has improved of late quite significantly with the improvement in the b -tagging efficiency. Its expression in the SM is given by

$$R_b^{\text{SM}} \simeq 0.220 + 0.25 \Delta V_b^t. \quad (11)$$

It may be noted that R_b is free from $\Delta\rho$ and α_S (hard bremsstrahlung).

The fourth observable is [14]

$$T = \frac{3}{59} \frac{\Gamma_{had}}{\Gamma_e} - \frac{30}{59} \gamma_e. \quad (12)$$

Within the SM

$$T^{\text{SM}} \simeq \frac{29}{59} + \frac{19}{59} \Delta V_b^t + \frac{\alpha_S}{\pi}. \quad (13)$$

In the theoretical expression of T , the $\Delta\rho$ term drops out.

Although ΔV_b^t contains a quadratic m_t dependence, it is free from M_H to a significant extent, and is also reasonably clean from the kinds of new physics we are looking at. On the other hand, $\Delta\rho$, in addition to having a quadratic m_t dependence, has a reasonable amount of M_H dependence (logarithmic); additionally, it receives tree-level contributions from the types of new physics under our consideration. So, although the presence of ΔV_b^t induces some amount of uncertainties through m_t , still the above $\Delta\rho$ -free observables can be used as ‘clean microscopes’ for looking at physics beyond the SM.

The present stage of experimental precision at LEP (after the analysis of the 1992 results) yields:

$$\begin{aligned} \tilde{\gamma}_e^{\text{exp}} &= 0.396 \pm 0.004 \\ R_{\tau e}^{\text{exp}} &= 0.996 \pm 0.007 \\ R_b^{\text{exp}} &= 0.2200 \pm 0.0027 \\ T^{\text{exp}} &= 0.513 \pm 0.005. \end{aligned} \quad (14)$$

We now intend to use these ‘microscopes’ to search for a new physics scenario comprising extra neutral gauge bosons and exotic fermions. We will examine each type of extension *one at a time* to prevent possible conspiring interplay among the different new parameters which will be present in the extended models. Effects of the simultaneous presence of more than one new physics will also be discussed in some cases.

3 Z^0 – $Z^{0'}$ mixing

Additional gauge bosons besides the ones predicted by the SM are present in many theoretically appealing extensions, which project out the SM as their low-energy manifestation. The Grand Unified Theories (GUTs) are potentially viable candidates for providing such neutral bosons. Many non-GUT models are also there with their own justification, such as the LRS model, which also accommodate additional gauge bosons. The additional gauge bosons can be neutral as well as charged. For the present purpose we consider the effects of extra $Z^{0'}$ bosons only as small perturbations to the SM. If, in particular, the gauge group is $SU(2) \otimes \prod_{\alpha=1}^n U(1)_\alpha$, there are $n-1$ extra $Z^{0'}$ bosons. Considering the fact that the couplings between the extra $U(1)$ bosons and the fermions are quite arbitrary, there are many unknown parameters, and the phenomenological analysis becomes cumbersome. But if we take $U(1)$'s as stemming from an underlying non-Abelian gauge group, in which there is only one overall coupling constant, the analysis becomes rather simple. The couplings of the extra $Z^{0'}$ bosons are then predicted from the GUT; the exact strengths are of course dependent on the symmetry breaking chains. When the gauge group $SU(2) \otimes U(1)^n$ breaks, at each stage of symmetry-breaking the corresponding intermediate vector bosons acquire masses. As a result of the diagonalization of the matrix consisting of the mass terms of the SM vector bosons, the masses of the extra gauge bosons and most importantly the off-diagonal terms that determine the strengths of mixings between the sequential and additional gauge bosons, the physical gauge boson states are different from the SM ones.

For the sake of simplicity we will consider the effects of only *one* extra $Z^{0'}$ and that of only two origins: (i) E_6 breaking down to the SM gauge group via intermediate breaking steps and yielding extra $U(1)$ group(s) (extra $U(1)$ models) and (ii) the LRS model with the gauge group $SU(2)_L \otimes SU(2)_R \otimes U(1)$ breaking in one step to the SM. In addition to taking the SM one-loop effects, we consider only the tree-level mixing effects of the additional $Z^{0'}$ with the SM Z^0 without making any assumption of the underlying Higgs structure responsible for the symmetry-breaking pattern.

The mass eigenstates (Z, Z') are obtained from the gauge eigenstates ($Z^0, Z^{0'}$) by making the following rotation

$$\begin{pmatrix} Z \\ Z' \end{pmatrix} = \begin{pmatrix} \cos \xi_0 & \sin \xi_0 \\ -\sin \xi_0 & \cos \xi_0 \end{pmatrix} \begin{pmatrix} Z^0 \\ Z^{0'} \end{pmatrix} \quad (15)$$

Apart from the introduction of this new mixing angle ξ_0 , the mass of the physical Z is changed from that of the SM Z^0 . The latter effect is manifested through a change of the tree-level ρ parameter. The modified vector (v_i) and axial vector (a_i) couplings of the physical Z to the standard fermions are given by

$$v_i = \sqrt{\rho} \left[t_3^i - 2Q_i \sin^2 \bar{\theta}_W + \xi_0 g'_{vi} \right]$$

$$a_i = \sqrt{\rho} [t_3^i + \xi_0 g'_{ai}] \quad (16)$$

where $\rho = 1 + \Delta\rho = 1 + \Delta\rho_{\text{tree}} + \Delta\rho_t = 1 + \Delta\rho_{\text{SB}} + \Delta\rho_{\text{M}} + \Delta\rho_t$ and $\sin^2 \bar{\theta}_W \simeq S_0^2 - \frac{3}{8}\Delta\rho$. The tree-order symmetry breaking (SB) and mixing (M) contributions, $\Delta\rho_{\text{SB}}$ and $\Delta\rho_{\text{M}}$ respectively, need not be evaluated separately, since $\Delta\rho$ does not appear explicitly in the expressions for the four observables chosen in this study. We examine only the mixing angles, that cast observable consequences on those observables.

For the extra $U(1)$ model

$$\begin{aligned} g'_{vu} &= 0 \\ g'_{au} &= \frac{2}{3} \sin \theta_W \cos \theta_2 \\ g'_{vd} &= \frac{1}{2} \sin \theta_W \left(\cos \theta_2 + \sqrt{\frac{5}{3}} \sin \theta_2 \right) \\ g'_{ad} &= -\sin \theta_W \left(\frac{-1}{6} \cos \theta_2 + \frac{1}{2} \sqrt{\frac{5}{3}} \sin \theta_2 \right) \\ g'_{ve} &= -\frac{1}{2} \sin \theta_W \left(\cos \theta_2 + \sqrt{\frac{5}{3}} \sin \theta_2 \right) \\ g'_{ae} &= g'_{ad} \\ g'_{v\nu} &= -\sin \theta_W \left(\frac{1}{6} \cos \theta_2 + \frac{1}{2} \sqrt{\frac{5}{3}} \sin \theta_2 \right) \\ g'_{a\nu} &= g'_{v\nu} \end{aligned} \quad (17)$$

where θ_2 is an angle characteristic of a given E_6 -symmetry-breaking chain yielding an extra $U(1)$. Four such models are of a general phenomenological interest corresponding to $\theta_2 = 0^\circ, 52.24^\circ, -52.24^\circ$ and -82.76° .

In the LRS model, with $\lambda = g_L/g_R$ and $y = \sqrt{\cos^2 \theta_W - \lambda^2 \sin^2 \theta_W}$,

$$\begin{aligned} g'_{vi} &= \frac{\cos^2 \theta_W}{\lambda y} (t_{3R}^i - 2\lambda^2 Q_i \sin^2 \theta_W) + \frac{\lambda \sin^2 \theta_W}{y} (t_{3L}^i - 2Q_i \sin^2 \theta_W) \\ g'_{ai} &= -\frac{\cos^2 \theta_W}{\lambda y} t_{3R}^i + \frac{\lambda \sin^2 \theta_W}{y} t_{3L}^i. \end{aligned} \quad (18)$$

For all practical purposes, one can put $\sin^2 \theta_W = S_0^2 = 0.23$ in eqs. (17) and (18).

4 Ordinary-exotic fermion mixing

When the bigger groups discussed above break down to the SM gauge group, in addition to yielding the extra gauge bosons, they also entail some additional fermions with

various $SU(2) \otimes U(1)$ representations. We classify all fermions as either ‘ordinary’ or ‘exotic’ according to their transformation properties under $SU(2)$. All the standard fermions, i.e. those contained in the SM, are called ‘sequential’ fermions. Their left-chiral fields transform as doublets and right-chiral fields as singlets under $SU(2)$. In addition, we consider three other kinds of non-sequential fermions. These are (i) vector singlets, which transform as singlets under $SU(2)$ both in the left- and right-handed sectors, (ii) vector doublets, which transform as doublets under $SU(2)$ both in the left- and right-handed sectors and (iii) mirror fermions, which transform as singlets in the left-handed sector and doublets in the right-handed ones. It may be noted that fermions transforming identically in the two sectors (left and right) do not have any axial couplings with the gauge bosons and are called ‘vector’ particles. To complete the nomenclature, we define all left-handed fermions occurring in doublets (irrespective of whether they are members of sequential or vector doublets) as ‘ordinary’ and all left-handed singlets (mirror families or vector singlets) as ‘exotic’. Similarly we define all right-handed singlets (sequential or vector singlets) as ‘ordinary’ and all right-handed doublets (mirror families or vector doublets) to be ‘exotic’. In general, one could speculate on larger varieties of *exotica*. But we have constrained ourselves to only the three types mentioned above, which follow from various theoretical extensions beyond the SM. Our analyses are, however, quite general as they correspond to a wide class of models without relying specifically on particular ones. These fermions are quite heavy (at least heavier than $M_Z/2$): otherwise they would have been produced in pairs at LEP. So the only way they can manifest themselves at LEP is through their mixing with their sequential counterparts having the same electric charge and colour quantum numbers. Fermions with exotic charges and colour assignments do not mix with the standard fermions as $U(1)_{\text{em}}$ and $SU(3)_c$ are unbroken. These types of fermions have, therefore, not been considered in our analysis.

To simplify our analysis we assume that exotic fermions mix with standard ones *diagonally*, i.e. one exotic fermion mixes with a unique standard flavour. This assumption automatically ensures the absence of tree level flavour-changing neutral currents (FCNCs) between light fermion generations, which have extremely tight experimental constraints. We also assume, to make life simpler, that there is only one exotic fermion at a time in the theory which mixes only with the *third*-generation flavour eigenstate. The loop effects of these fermions are also not considered.

In the presence of mixing the neutral current for the quarks and the charged leptons at tree level reads,

$$J_Z^\mu = \frac{g}{\cos\theta_W} \left[\bar{\psi}_{iL} \gamma^\mu \left(t_3^i C_L^{i2} - Q_i \sin^2 \theta_W \right) \psi_{iL} + \bar{\psi}_{iR} \gamma^\mu \left(t_3^i S_R^{i2} - Q_i \sin^2 \theta_W \right) \psi_{iR} \right], \quad (19)$$

where $C_L^i \equiv \cos \theta_L^i$ and $S_R^i \equiv \sin \theta_R^i$; θ_L^i and θ_R^i are interpreted as light–heavy mixing angles in the left- and right-handed sectors respectively. The C_L^{i2} term represents a non-universal reduction of strength of the normal neutral current due to mixing with

left-handed singlets, and the $S_R^i{}^2$ term represents an induced right-handed current, which is generated as a result of mixing with right-handed doublets.

The effective vector and axial vector couplings of the Z to the fermion i follow immediately as

$$\begin{aligned} v_i &= \sqrt{\rho} \left[t_3^i (C_L^i{}^2 + S_R^i{}^2) - 2Q_i \sin^2 \bar{\theta}_W \right] \\ a_i &= \sqrt{\rho} t_3^i (C_L^i{}^2 - S_R^i{}^2). \end{aligned} \quad (20)$$

It may be noted that the SM couplings (given in eq. (3)) can be recovered from eq. (20) by setting $\theta_L = \theta_R = 0$. Note that such mixing angles have observable consequences only when there are mixings between states with different t_3 . For example, when a b quark mixes with a singlet h (charge $-1/3$ non-sequential, which stems from E_6), only θ_L is non-trivial while when b mixes with a B quark sitting in a vector doublet, only θ_R is relevant. Similar arguments hold for the leptonic sector as well. We do not treat the mirror fermions separately as their mixing effects will always mimic the joint impact of mixings in the left- and right-handed sectors that we consider case by case in this analysis.

The mixing effects of the neutrinos follow the same textures as those of the charged fermions. But the framework is a bit more complicated due to the presence of the Majorana mass terms of the neutrinos and due to the lack of experimental constraints on neutrino FCNCs. In the present analysis we do not deal with neutrino mixings.

5 Results

The reduced electronic width can provide the cleanest (because ordinary–exotic fermion mixings are negligibly small in the first two generations) and the most severe bounds on the possible Z^0 – $Z^{0'}$ mixing effects. Such kinds of mixings modify the reduced electronic width as (using first-order approximation)

$$\begin{aligned} \tilde{\gamma}_e &\simeq \tilde{\gamma}_e^{\text{SM}} + 4 \sin \theta_W \left(-\frac{1}{6} \cos \theta_2 + \frac{1}{2} \sqrt{\frac{5}{3}} \sin \theta_2 \right) \xi_0 \quad (\text{extra U(1)}) \\ \tilde{\gamma}_e &\simeq \tilde{\gamma}_e^{\text{SM}} - 2 \sqrt{\cos 2\theta_W} \xi_0 \quad (\text{LRS}). \end{aligned} \quad (21)$$

The bounds are displayed in Table 1. It is seen that for different models under consideration, $|\xi_0|$ lies in the range of (1–5)% at 95% confidence level (C.L.). These bounds are stronger than the ones in the previous analyses, owing to the reduction of the systematic and statistical errors of the LEP measurements.

If the τ lepton mixes with a vector singlet (or with its exotic partner sitting in a vector doublet), then neglecting the Z^0 – $Z^{0'}$ mixing, one obtains,

$$\begin{aligned} \tilde{\gamma}_\tau &\simeq \tilde{\gamma}_e^{\text{SM}} - 2S_L^{\tau 2} \quad (\text{singlet}) \\ \tilde{\gamma}_\tau &\simeq \tilde{\gamma}_e^{\text{SM}} - 2S_R^{\tau 2} \quad (\text{doublet}) \end{aligned} \quad (22)$$

When the simultaneous presence of τ -exotic (τ) and Z^0 - $Z^{0'}$ mixing is considered, the best observable to put a bound on the former, keeping it free from the uncertainties of the latter (employing the generation universality of $Z^{0'}$ coupling to leptons), is $R_{\tau e}$; it is given by

$$\begin{aligned} R_{\tau e} &\simeq 1 - 2S_L^{\tau^2} \quad (\text{singlet}) \\ R_{\tau e} &\simeq 1 - 2S_R^{\tau^2} \quad (\text{doublet}). \end{aligned} \quad (23)$$

for mixing of τ with a vector singlet (vector doublet). The bounds at 95% C.L. are shown in Table 2.

When the b quark mixes with a vector singlet (or a member of a vector doublet) partner, and assuming there is no Z^0 - $Z^{0'}$ mixing,

$$\begin{aligned} R_b &\simeq R_b^{\text{SM}} - 0.40 S_L^{b^2} \quad (\text{singlet}) \\ R_b &\simeq R_b^{\text{SM}} - 0.08 S_R^{b^2} \quad (\text{doublet}). \end{aligned} \quad (24)$$

The bounds at 95% C.L. of the experimental uncertainties are shown in Table 2. It may be noted that the constraints on $S_L^{b^2}$ and $S_R^{b^2}$ are significantly stringent at $m_t = 200$ GeV compared to the ones for $m_t = 100$ GeV, because the experimental central value of R_b at present is close the theoretical prediction for smaller values of m_t . Also to be noted is that the effect of mixing with a doublet is 5 times less sensitive than with a singlet.

If the Z^0 - $Z^{0'}$ mixing (of extra $U(1)$ type, say) and b - h (say) mixings are considered simultaneously, then (using first order approximation)

$$R_b \simeq R_b^{\text{SM}} - 0.40 S_L^{b^2} + \alpha(\theta_2) \xi_0 \quad (25)$$

where

$$\alpha(\theta_2) \simeq \sin \theta_W (-0.28 \cos \theta_2 + 0.04 \sin \theta_2). \quad (26)$$

Just to feel the numerical impact of such simultaneous mixing we take, as an example, the $\theta_2 = 0^\circ$ model, $m_t = 100$ GeV, and employ the bounds on $|\xi_0|$ from $R_{\tau e}$, to obtain $S_L^{b^2} \leq 0.028$ at 95% C.L. As a result the bound is seen to have been relaxed when one compares it with the corresponding one, namely $S_L^{b^2} \leq 0.011$ (see Table 2), in the absence of Z^0 - $Z^{0'}$ mixing.

The expression of the T parameter in the presence of mixing of a b quark with a singlet becomes

$$T \simeq T^{\text{SM}} - \frac{30}{59} S_L^{b^2}. \quad (27)$$

The corresponding bounds at 95% C.L. are shown in Table 2. We choose $\alpha_S = 0.12$. Uncertainties due to α_S are small. Mixing with a vector doublet, as has been seen in the context of R_b , relaxes the bound by a factor of 5. Thus all the bounds on ordinary-exotic fermion mixing angles lie in the range of (1-5)% at 95% C.L. except for mixing with vector doublets, where the mixing angles could be large.

6 Conclusion

To conclude, we have examined the Z^0 – $Z^{0'}$ and ordinary–exotic mixings in the light of the precision data obtained from ~ 5 million Z events. These bounds have been derived from variables that do not depend on $\Delta\rho$. Most of these mixing angles, including all Z^0 – $Z^{0'}$ ones, are found to be almost vanishing. For these cases it may be argued that to keep the mass(es) of Z' or of the corresponding exotic fermions in the accessible range of the forthcoming colliders, favourable choices of the underlying Higgs structure are necessary; these require closer scrutiny. On the other hand, some mixing angles (particularly the right-handed ones for b quark) could still be significantly large. Further reduction of the systematic and statistical errors of the LEP measurements and/or the discovery of the top quark would definitely reheat these issues in attempts to search for exotic fermions through such indirect probes.

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Table 1: Upper bounds on Z^0 – $Z^{0'}$ mixing angles at 95% C.L. The observable from which the bounds are derived is $\tilde{\gamma}_e$.

Extra $U(1)$					LRS
θ_2 (deg)	0	52.24	−52.24	−82.76	$\lambda = 1$
$ \xi_0 $ (rad)	0.05	0.02	0.02	0.01	0.01

Table 2: Upper bounds on ordinary–exotic mixing angles at 95% C.L. The observables from which the bounds are derived are mentioned. For b quark mixing angles the bounds refer to $m_t = 100$ (200) GeV, and they correspond to the situation when there is no Z^0 – $Z^{0'}$ mixing.

$R_{\tau e}$		R_b		T	
$S_L^{\tau^2}$	$S_R^{\tau^2}$	$S_L^{b^2}$	$S_R^{b^2}$	$S_L^{b^2}$	$S_R^{b^2}$
0.009	0.009	0.011 (0.001)	0.055 (0.005)	0.05 (0.04)	0.25 (0.20)